

Heavy Hadron Spectrum and Interactions

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Starting from the approximate symmetries of QCD, namely chiral symmetry for light quarks and spin and flavor symmetry for heavy quarks, we investigate the low-energy properties of heavy hadrons. For this purpose we construct a consistent picture of quark-antiquark and quark-diquark interactions as a low-energy approximation to the flavor dynamics in heavy mesons and heavy baryons, respectively. Using standard functional integration tools, we derive an effective Lagrangian in terms of heavy hadron fields and discuss several properties, like the mass spectrum, coupling and decay constants, Isgur-Wise form factors.

1 Motivation

In the framework of the Standard Model, the theory of strong interaction is given by Quantum Chromodynamics (QCD), which is formulated as a renormalizable non-abelian gauge theory in terms of quark and gluon fields. Due to asymptotic freedom of the theory, the short distance part of the QCD dynamics can be treated perturbatively. However from the phenomenological point of view, one would like to formulate the theory in terms of the one-particle states entering the detectors in high-energy experiments, which are the hadrons rather than quarks and gluons. The long-distance QCD effects leading to the confinement of quark and gluons and their binding to hadrons, can only be understood non-perturbatively, and therefore it is a big challenge to derive effective hadron lagrangians directly from QCD. Such a program evidently requires (possibly crude) approximations on the long way from QCD down to the effective hadron theory.

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Simplifications arise, however, since we are naturally restricted to the low-energy region. Therefore, from the practical point of view, it is sufficient to find an approximation to QCD which mimics the essential features of low-energy quark flavor dynamics. Clearly, the possible form of such effective quark lagrangians must be restricted by the underlying symmetries of QCD, which should be viewed as a guide to find tractable models of quark flavor dynamics.

In the sector of light quark flavors $q = (u, d, s)$, QCD possesses an approximate $SU(3)_L \times SU(3)_R$ chiral symmetry which is spontaneously broken to $SU(3)_V$, leading to the emergence of (pseudo)Goldstone bosons π, K, η , which receive their masses by the explicit breaking of chiral symmetry through current quark masses.

Recently, new important symmetries have also been discovered for heavy quark flavors $Q = b, c, \dots$ which considerably simplify the description of hadrons containing one heavy quark. These symmetries arise in the limit of infinite heavy quark masses $m_Q \rightarrow \infty$, where the heavy quark spin decouples from the light QCD degrees of freedom, and the dynamics of the light quarks within a heavy hadron is independent of the heavy quark flavor. Consequently, heavy hadrons are organized in spin symmetry multiplets, and flavor symmetry tells us how heavy hadron observables scale with the heavy mass. Corrections to the heavy quark limit are treated systematically in the Heavy Quark Effective Theory (HQET) [1, 2, 3].

For light quark flavors the Nambu–Jona–Lasinio (NJL) model has been successfully used to describe the dynamics of QCD. The properties of this model are governed by global chiral invariance and its spontaneous breaking in the ground state. Indeed, employing appropriate functional integration techniques the model can be reformulated as an effective low-energy lagrangian in terms of light pseudoscalar, vector and axial vector mesons fields, which are described surprisingly well, embodying the soft-pion theorems, vector dominance, Goldberger–Treiman and KSFR relations and the integrated chiral anomaly [4]. For a recent review on these subjects see ref. [5] and references therein.

In this talk we shall review some recent work on heavy meson physics based on an extension of the NJL model [6] which includes chiral symmetry for light quarks and the heavy quark symmetries (for related work see also [7]). In case of baryons, the diquark concept has proven extremely useful to approximate the dynamics at low energies. We will present our results based on a phenomenological ansatz for an interaction between heavy quarks and light diquarks [8] and a Faddeev-equation involving composite light and heavy diquarks [9], respectively.

2 Basic features of HQET

In order to extract the heavy quark symmetries from QCD one first identifies the relevant degrees of freedom for the heavy quark field $Q(x)$. These are given by the ‘upper component’

$$Q(x) \rightarrow Q_v(x) = \frac{1 + \not{v}}{2} e^{i m_Q v x} Q(x) . \quad (1)$$

Here the dependence on the heavy quark mass has been made explicit by means of the heavy quark velocity v_μ with $v^2 = 1$. Upon integrating out the irrelevant degrees of freedom, the (tree-level) HQET Lagrangian can be obtained as a $1/m_Q$ series of local operators [3]

$$\mathcal{L}_0^{\text{HQET}} = \bar{Q}_v (iv \cdot \mathcal{D}) Q_v + O(1/m_Q) , \quad (2)$$

where \mathcal{D}^μ is the covariant derivative in QCD. The operator $(iv \cdot \mathcal{D})$ obviously does not depend on the heavy quark spin and flavor, and therefore in the limit $m_Q \rightarrow \infty$ we discover the heavy quark spin and flavor symmetries.

One of the striking consequences of these symmetries can be formulated as an analogue of the Wigner-Eckhart-theorem. Using a matrix representation for the heavy pseudoscalar and vector meson fields (describing the $B - B^*$ and $D - D^*$, respectively²)

$$\mathcal{H}_v = \sqrt{M_H} \frac{1 + \not{v}}{2} [i\gamma_5 + \not{\epsilon}]$$

the hadronic matrix elements of any heavy quark current can be expressed through

$$\langle \bar{H}(v) | \bar{Q}_{v'} \Gamma Q_v | H(v) \rangle = \xi(v \cdot v') \text{tr} [\bar{\mathcal{H}}_{v'} \Gamma \mathcal{H}_v] . \quad (3)$$

Here the universal Isgur-Wise formfactor [2] $\xi(v \cdot v')$ plays the role of a reduced matrix element, and the Dirac trace can be viewed as a Clebsch-Gordan coefficient. Furthermore for $v = v'$ the normalization of the Isgur-Wise function is restricted by the symmetries to be unity $\xi(1) = 1$. Similar statements are true for heavy baryons (see below).

3 Heavy Mesons

Let us consider an extended NJL model containing besides the free lagrangian of light quarks (q) and heavy quarks (Q_v) of definite velocity v , a four-quark interaction term which is motivated by the general quark-current structure of QCD and obeys chiral and heavy quark symmetries. After Fierz-rearrangement into the physical color-singlet channel the relevant interaction term between a light quark and a heavy quark is given by (summation over repeated indices is understood) [6]

$$\begin{aligned} \mathcal{L}_{int}^{hl} = & G_3 ((\bar{Q}_v i\gamma_5 q)(\bar{q} i\gamma_5 Q_v) - (\bar{Q}_v \gamma^\mu q) P_{\mu\nu}^\perp (\bar{q} \gamma^\nu Q_v)) \\ & + G_3 ((\bar{Q}_v q)(\bar{q} Q_v) - (\bar{Q}_v i\gamma^\mu \gamma_5 q) P_{\mu\nu}^\perp (\bar{q} i\gamma_5 \gamma^\nu Q_v)) , \end{aligned} \quad (4)$$

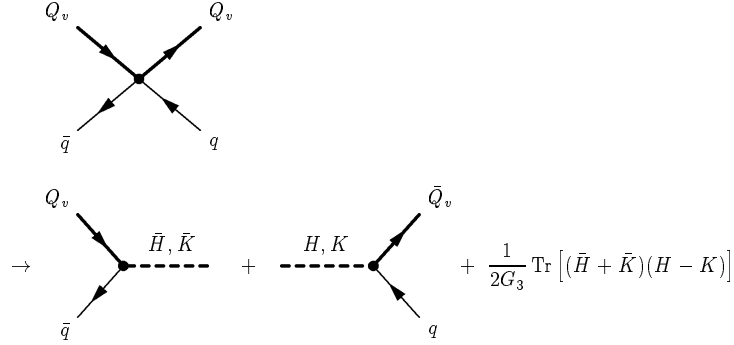
where $P_{\mu\nu}^\perp = g_{\mu\nu} - v_\mu v_\nu$ is a projection operator which is transversal with respect to the heavy quark velocity, $v^\mu P_{\mu\nu} = 0$.

The bosonization procedure is standard [4, 5] and consists in introducing composite mesonic fields $(\bar{q}q)$, $(\bar{q}Q_v)$ into the generating functional which is given by the path integral

$$Z = \int \mathcal{D}q \mathcal{D}\bar{q} \mathcal{D}Q_v \mathcal{D}\bar{Q}_v e^{i \int d^4x (\mathcal{L}_0 + \mathcal{L}_{int}^{ll} + \mathcal{L}_{int}^{hl})} , \quad (5)$$

²Here ϵ_μ denotes the polarization vector of the vector mesons with $v \cdot \epsilon = 0$.

such that the lagrangian becomes bilinear in the quark fields and the latter can be integrated out. Here \mathcal{L}_{int}^{ll} is the usual interaction term of light quarks. Symbolically, one has



We use a non-linear representation where $\xi = \exp(i\pi/F)$ is an element in the coset space $SU(3)_L \times SU(3)_R / SU(3)_V$. Here F is the bare decay constant, and $\pi = \pi^a \lambda_F^a / 2$ represents the light octet of (pseudo)Goldstone bosons associated to spontaneous breakdown of chiral symmetry. Our model includes light vector mesons $V_\mu = V_\mu^a \lambda_F^a / 2$ and axial-vector mesons $A_\mu = A_\mu^a \lambda_F^a / 2$.

In the heavy-light sector we can collect the pseudoscalar field Φ^5 and the vector field Φ^μ into a (super)field H which represents the $J^P = (0^-, 1^-)$ doublet of spin symmetry. Analogously the scalar field Φ and the axial-vector field $\Phi^{5\mu}$ are combined in the (super)field K describing the $J^P = (0^+, 1^+)$ doublet

$$H = \frac{1 + \not{v}}{2} (i\Phi^5 \gamma_5 + \Phi^\mu \gamma_\mu) \quad , \quad v_\mu \Phi^\mu = 0 \quad , \quad (6)$$

$$K = \frac{1 + \not{v}}{2} (\Phi + i\Phi^{5\mu} \gamma_\mu \gamma_5) \quad , \quad v_\mu \Phi^{5\mu} = 0 \quad . \quad (7)$$

Due to flavor symmetry of HQET these fields describe both B or D mesons. The details can be found in ref. [6].

Performing the integration over heavy and light quark fields in (5) gives the quark determinant contributing to the effective lagrangian as follows

$$\mathcal{L}_{\text{eff}} = -iN_c \text{Tr} \ln i\mathcal{D} + \frac{1}{2G_3} \text{Tr} [(\overline{H} + \overline{K})(H - K)] + \dots \quad , \quad (8)$$

where

$$i\mathcal{D} = i\cancel{\partial} - \Sigma + \not{V} + \not{A}\gamma_5 - (\overline{H} + \overline{K})(iv \cdot \partial)^{-1}(H + K) \quad (9)$$

is the quark Dirac operator containing light and heavy meson fields.

To regularize the quark loops arising from (8) we shall use a universal proper-time cut-off Λ which will be fixed from the light meson data. For the applicability of the

NJL-model to heavy quark dynamics it is crucial to observe that the typical relative momentum in a heavy hadron is set by the light degrees of freedom (gluons, light quarks) only, since all reference to the heavy quark mass has been removed through (1) and (2). Therefore it is the residual momentum of the heavy quark $k_\mu = P_\mu - m_Q v_\mu$ which is regularized by the cut-off Λ .

3.1 Heavy meson self-energy and weak decay constant

Expanding the term $-iN_c \text{Tr} \ln i\bar{D}$ in (8) in powers of the meson fields leads to the familiar loop expansion given by Feynman diagrams with heavy and light mesons as external lines and heavy and light quarks in internal loops. For the light sector this has been done to derive an effective lagrangian in terms of π , ρ and A_1 fields [4]. Comparison with experimental data fixes the parameters relevant for the heavy sector: the light constituent quark masses, $m^{u,d} = 300$ MeV, $m^s = 510$ MeV, and a universal cut-off $\Lambda = 1.25$ GeV.

The self-energy part for the heavy mesons can be expanded in powers of the external momentum $v \cdot p$, such that the effective meson lagrangian acquires in configuration space the desired form

$$\mathcal{L}_0^H = -\text{Tr} [\bar{H} (iv \cdot \partial - \Delta M_H) H] + \text{Tr} [\bar{K} (iv \cdot \partial - \Delta M_K) K] \quad , \quad (10)$$

where the mass differences between heavy meson and heavy quarks are $\Delta M_{H,K} = M_{H,K} - m_Q$. Any left-handed weak decay current $\bar{q}_L \Gamma Q_v$ of a heavy meson H can be represented by

$$J_\Gamma = \frac{\sqrt{M_H} f_H}{2} \text{Tr} [\xi^\dagger \Gamma H]$$

with f_H being the weak decay constant defined through

$$\langle 0 | \bar{q} \gamma_\mu (1 - \gamma_5) Q_v | H_v(0^-) \rangle = i f_H M_H v_\mu \quad .$$

In the extended NJL model it can be expressed in a simple way by the Z -factors of meson fields and the four-quark coupling constant G_3

$$f_H \sqrt{M_H} = \sqrt{Z_H} / G_3 \quad . \quad (11)$$

We thus recover the familiar scaling of the weak decay constant of heavy mesons with the heavy mass in HQET due to heavy spin and flavor symmetry. A similar relation holds also for the members of the $J^P = (0^+, 1^+)$ doublet. The experimental result for masses and decay constants as a function of the four-quark coupling constant G_3 is presented in Table 1.

3.2 Strong interactions with Goldstone bosons

The Goldstone bosons π , K , η couple to heavy meson fields via the vector and axial-vector combinations

$$\mathcal{V}_\mu^\pi = \frac{i}{2} (\xi \partial_\mu \xi^\dagger + \xi^\dagger \partial_\mu \xi) \quad , \quad \mathcal{A}_\mu^\pi = \frac{i}{2} (\xi \partial_\mu \xi^\dagger - \xi^\dagger \partial_\mu \xi) \quad .$$

Table 1: Heavy meson parameters as a function of the coupling constant G_3 .

$G_3[\text{GeV}^{-2}]$	3	5	7	9	11	expected
$\Delta M_H^s - \Delta M_H^{u,d} [\text{MeV}]$	250	160	110	80	50	100
$\Delta M_K^{u,d} - \Delta M_H^{u,d} [\text{MeV}]$	370	390	410	430	450	?
$f_B [\text{MeV}]$	300	210	170	150	130	150–180
f_H^s/f_H^u	1.13	1.13	1.12	1.11	1.09	1.1–1.2

The vector field \mathcal{V}_μ^π essentially couples with a covariant derivative of $SU(3)_V$

$$\partial_\mu H \rightarrow D_\mu H = \partial_\mu H + i H \mathcal{V}_\mu^\pi$$

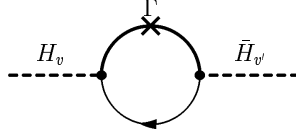
with small corrections due to the explicit breaking $m_s \neq m_{u,d}$. The coupling of the axial vector field

$$g_{HHA} \text{tr}_D [\bar{H} H \not{A}^\pi \gamma_5]$$

is not restricted by the symmetries. Our model predicts a value of about $g_{HHA} \approx -0.2$ in accordance with experimental bounds $\Gamma(D^{*+} \rightarrow D^0 \pi^+) < 0.131 \text{ MeV} \Rightarrow g_{HHA}^2 < 0.5$.

3.3 Isgur-Wise function

The calculation of the Isgur-Wise function defined in eq. (3) involves the following Feynman diagram



As a result we obtain for external momentum $v \cdot p = 0$ the general form

$$\xi(\omega) = a \frac{2}{1+\omega} + (1-a) \frac{\ln(\omega + \sqrt{\omega^2 - 1})}{\sqrt{\omega^2 - 1}}, \quad a = \text{const.} \quad (12)$$

Numerically for the slope parameter of the Isgur-Wise function at the non-recoil point $\rho^2 = -\xi'(1)$ we obtain the estimate $\rho = 0.67$. Other recent estimates are slightly larger with $0.84 \leq \rho \leq 1.0$ (see [10] and references therein).

4 Heavy baryons

Light baryons have been recently described within quark models as bound states of quarks and diquarks [5]. In ref. [8] we have performed a first step to obtain an analogous description of heavy baryons (containing one heavy quark) in the heavy quark limit by making use of the diquark picture. Light diquarks are then coupled to the heavy quark

by a simple model interaction. The lowest-lying baryons containing a single heavy quark are classified according to their flavor content (see Table 2). The light quark quantum numbers (color, flavor, spin) within a diquark are coupled as follows

$$3_C \otimes 3_C = \bar{3}_C \oplus 6_C, \quad 3_F \otimes 3_F = \bar{3}_F \oplus 6_F, \quad (1/2) \otimes (1/2) = 0 \oplus 1. \quad (13)$$

Table 2: Flavor content of the lowest-lying heavy baryons.

charmed baryon	spin partner	bottom baryon	notation	ij	I	S
Λ_c^+	trivial	Λ_b^0	T_v^{ij}	$[ud]$	0	0
$\Xi_c^{+,0}$	trivial	$\Xi_b^{0,-}$	T_v^{ij}	$[us], [ds]$	1/2	-1
$\Sigma_c^{++,+,0}$	$\Sigma_c^{*+++,+0}$	$\Sigma_b^{(*)+,0,-}$	$S_v^{\mu ij}$	$uu, \{ud\}, dd$	1	0
$\Xi_c'^{+,0}$	$\Xi_c^{*+,0}$	$\Xi_b'^{(*)0,-}$	$S_v^{\mu ij}$	$\{su\}, \{ds\}$	1/2	-1
Ω_c^0	Ω_c^{*0}	$\Omega_b^{(*)-}$	$S_v^{\mu ij}$	ss	0	-2

Clearly, one has to consider only the color antitriplet part in (13) which together with the color triplet of the heavy quark forms the physical color singlet baryon state. The Pauli principle then requires the light diquark system to be symmetric with respect to simultaneous exchange of spin and flavor indices. Therefore scalar diquarks only occur in the flavor antitriplet $\bar{3}_F$ and are hence realized by an antisymmetric flavor matrix D^{ij} , while axial-vector diquarks are realized by the symmetric 6_F flavor matrix F^{ij} .

Analogously, in the heavy mass limit we have two types of heavy baryons: In the first case the two light quarks are coupled to spin zero, and the heavy baryons are represented by an anti-symmetric flavor matrix of ordinary spin-1/2 Dirac spinors

$$T_v^{ij} = \frac{i}{\sqrt{2}} \begin{pmatrix} 0 & \Lambda_Q & \Xi_{Q, I_3=1/2} \\ -\Lambda_Q & 0 & \Xi_{Q, I_3=-1/2} \\ -\Xi_{Q, I_3=1/2} & -\Xi_{Q, I_3=-1/2} & 0 \end{pmatrix}. \quad (14)$$

The two spin orientations of T_v form the (in this case trivial) spin symmetry partners. In the second case the two light quarks are coupled to spin one and one obtains a symmetric flavor matrix of either spin-1/2 Dirac fields B_v^{ij} or spin-3/2 Rarita-Schwinger fields $B_v^{*\mu ij}$. These are then symmetry partners with respect to spin rotations of the heavy quark. Consequently, they can be combined in a multiplet

$$S_v^{\mu ij} = \frac{1}{\sqrt{3}} \gamma_5 (\gamma^\mu - v^\mu) B_v^{ij} + B_v^{*\mu ij}, \quad v_\mu S_v^\mu = 0$$

$$B_v^{(*)ij} = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2} \Sigma_{Q, I_3=1}^{(*)} & \Sigma_{Q, I_3=0}^{(*)} & \Xi_{Q, I_3=1/2}'^{(*)} \\ \Sigma_{Q, I_3=0}^{(*)} & \sqrt{2} \Sigma_{Q, I_3=-1}^{(*)} & \Xi_{Q, I_3=-1/2}'^{(*)} \\ \Xi_{Q, I_3=1/2}'^{(*)} & \Xi_{Q, I_3=-1/2}'^{(*)} & \sqrt{2} \Omega_Q^{(*)} \end{pmatrix}. \quad (15)$$

4.1 Heavy quarks coupled to light diquarks

A straightforward phenomenological ansatz for the interaction of heavy and light quarks in the heavy baryon consists in treating the two light quarks as elementary scalar and

axial vector diquark fields D^{ij} , F_μ^{ij} , respectively. The interaction term which respects chiral and heavy quark symmetries is expressed as [8]

$$\mathcal{L}_{\text{int}} = \tilde{G}_1 \text{tr} [\bar{Q}_v D^\dagger D Q_v] - \tilde{G}_2 \text{tr} \left[\bar{Q}_v F_\mu^\dagger \left([1/2] P^{\mu\nu} + [3/2] P^{\mu\nu} \right) F_\nu Q_v \right] , \quad (16)$$

where \tilde{G}_1 , \tilde{G}_2 are effective coupling constant of dimension mass⁽⁻¹⁾ and

$$[1/2] P^{\mu\nu} = \frac{1}{3} (\gamma^\mu - \not{v} v^\mu) (\gamma^\nu - \not{v} v^\nu) , \quad [3/2] P^{\mu\nu} = g^{\mu\nu} - v^\mu v^\nu - [1/2] P^{\mu\nu} ,$$

$$v_\mu [1/2] P^{\mu\nu} = v_\mu [3/2] P^{\mu\nu} = \gamma_\mu [3/2] P^{\mu\nu} = 0 .$$

Again, we introduce the heavy baryon fields T^{ij} and S_μ^{ij} directly in the functional integral such that the quark and diquark fields can be integrated out. This leads to a quark-diquark loop expansion of the resulting functional determinant (for details see [8]). The self-energy part for the heavy baryon fields reads

$$\mathcal{L}_0^{\text{eff}} = \text{tr}_F \bar{T}_v (iv \cdot \partial - \Delta M_T) T_v - \text{tr}_F \bar{S}_{v\mu} (iv \cdot \partial - \Delta M_S) S_v^\mu , \quad (17)$$

where the residual masses $\Delta M_{T(S)} = M_{T(S)} - m_Q$ are related to the coupling parameters \tilde{G}_1 (\tilde{G}_2), the diquark masses $M_{D(F)}$ and the cut-off parameter used to regularize the quark-diquark loops.

An interesting application of this simple model is the calculation of the Isgur-Wise form factors for weak heavy baryon transitions, which are defined by the following expressions [11]

$$A(v \cdot v') \bar{T}_v \Gamma T_{v'} , \quad \{ B(v \cdot v') g_{\mu\nu} + C(v \cdot v') v'_\mu v_\nu \} \bar{S}_v^\mu \Gamma S_{v'}^\nu . \quad (18)$$

The Isgur-Wise formfactors $A(\omega)$ and $B(\omega)$ are normalized at zero recoil $A(1) = B(1) = 1$. Inserting an arbitrary weak current into the heavy quark line and calculating the one-loop Feynman diagrams, we obtain

$$A(\omega) = r(\omega) = \frac{\ln(\omega + \sqrt{\omega^2 - 1})}{\sqrt{\omega^2 - 1}} ,$$

$$B(\omega) = r(\omega) + \kappa \frac{\omega - r(\omega)}{2 + \omega^2} , \quad C(\omega) = -\kappa \frac{2 + \omega r(\omega)}{2 + \omega^2} , \quad (19)$$

with $\kappa > 0$ depending on the model parameters. Although the parameter range of our model cannot be restricted sufficiently by experimental data it is worth comparing our results with the theoretical constraints imposed by Bjorken and Xu [12]. Following the discussion in ref. [13], one obtains for $\omega \geq 1$

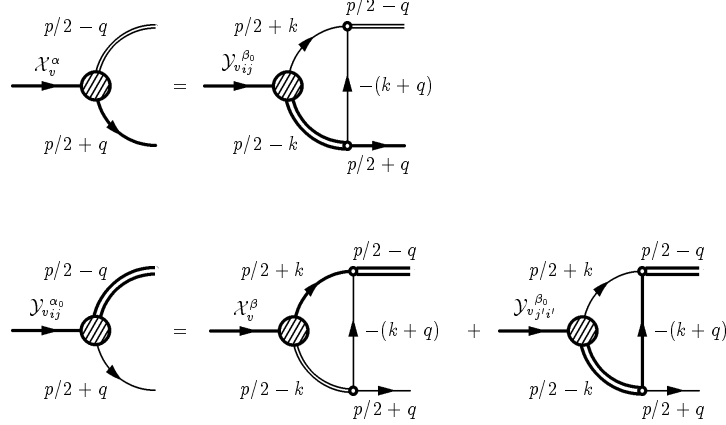
$$A(\omega) \leq 1 , \quad A'(1) \leq 0 ,$$

$$\frac{2}{3} |B(\omega)|^2 + \frac{1}{3} |\omega B(\omega) + (\omega^2 - 1) C(\omega)|^2 \leq 1 , \quad B'(1) \leq -\frac{1}{3} - \frac{2}{3} C(1) , \quad (20)$$

which is true in our model for any positive value of κ .

4.2 Light and heavy diquarks in the heavy baryon

In a more sophisticated description of heavy baryons we consider light **and** heavy diquarks as **composite** fields, with their interaction given by the attractive diquark channel included in the extended NJL model (see [9] for technical details). Upon introducing composite diquark and baryon fields with suitable constraints in the functional integral and performing several functional integrations, we obtain a system of coupled equations (Faddeev equation) for the two configurations $\mathcal{X} \sim Q [q\bar{q}^c]$ and $\mathcal{Y} \sim q [Q\bar{q}^c]$ describing the heavy baryon mass spectrum. In our context the interaction takes place via quark exchange, as it is illustrated below



Using a static approximation to the quark exchange and fixing some model parameters from the light baryon sector, we obtain a rough but reasonable estimate of heavy baryon masses (see Table 3).

Table 3: Heavy baryon mass splittings in MeV. The NJL coupling constant G_1 , which triggers the diquark properties, is given in units of the coupling constant of the (light) pseudoscalar meson sector.

G_1/G_{meson}	$\Sigma_Q - \Lambda_Q$	$\Xi_Q - \Lambda_Q$	$\Omega_Q - \Lambda_Q$	$\Xi_Q^{(*)} - \Xi_Q$
1.20	160	190	490	135
1.30	175	200	515	145
1.40	190	200	530	155
(Exp./Latt.)	170 [14]	185 [14]	455 [15]	155 [16]

5 Conclusions

In the present talk we have shown how chiral and heavy quark symmetries can be implemented in relativistic quark and quark-diquark models in order to describe the properties of heavy hadrons. We have investigated the extended NJL model, including

$SU(3)_F$ breaking and discussed the resulting mass spectrum for heavy mesons of the spin/parity doublets $J^P = (0^-, 1^-)$ and $J^P = (0^+, 1^+)$, their strong coupling constants and weak decay constants and the Isgur-Wise form factors.

Heavy baryons can be reasonably described within a model where the heavy quarks couple to elementary light scalar and axial vector diquark fields. The predicted form of the Isgur-Wise functions obeys the Bjorken/Xu inequalities in a parameter-independent way. On the basis of composite light and heavy diquark fields within the heavy baryon, a Faddeev equation is derived from the extended NJL model, which gives a reasonable description of the heavy baryon spectrum.

In order to improve the results on interactions, decays, Isgur-Wise functions etc., their remain several tasks to be done. One important question is how to settle the lack of confinement in NJL-type models which influences the external momentum-dependence of hadron properties. Another interesting problem is to what extent $1/m_Q$ corrections can be studied in such a framework.

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